

Metastable Rank-Condition Supersymmetry Breaking in a Chiral Example

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Abstract

We discuss generalizations of Intriligator-Seiberg-Shih (ISS) vacua to chiral models. We study one example, of an s-confining theory, in detail. In the IR, this example reduces to two ISS-like sectors, and exhibits a supersymmetry-breaking vacuum with all pseudo-moduli stabilized at the origin, and with the R-symmetry unbroken. The IR theory is interesting from the point of view of R-symmetry breaking. This theory is an O’Raifeartaigh model with all charges zero or two, but the presence of a second R-charged pseudo-modulus with superpotential couplings to the messengers in principle allows for R-symmetry breaking.

I. INTRODUCTION

Many examples of dynamical supersymmetry breaking theories are known, but these examples are rather non-generic [1]. If one gives up however on the requirement of a global supersymmetry-breaking vacuum, which is not essential for model building purposes anyway, the catalog of dynamical supersymmetry breaking theories may be greatly expanded. The discovery by Intriligator, Seiberg, and Shih (ISS) that metastable, supersymmetry-breaking vacua exist in simple theories such as SQCD [2], suggests that local dynamical supersymmetry breaking may indeed be a generic phenomenon. In this note, we take a first step towards searching for ISS-like minima in chiral models. To the best of our knowledge, no such examples are currently known, although there is no reason of principle that chiral theories could not exhibit local supersymmetry-breaking vacua.

In order to obtain calculable minima in which supersymmetry is dynamically broken via rank-conditions, one must find chiral theories whose IR descriptions involve some tensor (under the global symmetry) fields with dynamically-generated cubic superpotential couplings. s-confining theories [3, 4] provide a good starting point for this search, since they have smooth moduli-spaces and their IR behavior is well understood. Of the ten classes of s-confining $SU(N)$ theories [4], seven are chiral, in the sense that one cannot give mass to all the fields of the theory. Many of these have fields in 2-index (or higher) representations of the global symmetry group in the IR, but these fields often have only non-renormalizable superpotential couplings. In fact, only two models allow for rank-condition breaking: one is an $SU(5)$ theory with two copies of $(10 + \bar{5})$ plus two extra flavors [5], and the other is an $SU(6)$ theory with a 3-index anti-symmetric tensor and 4 flavors. Here we will focus on the latter because of its simplicity. Admittedly, this simplicity is related to the fact that the anti-symmetric tensor is a pseudo-real representation of $SU(6)$, so that the IR theory is similar to a vector-like theory. In fact, the IR theory essentially consists of two sets of ISS SQCD, plus one gauge invariant that couples to these fields only through non-renormalizable terms. Still, the theory is chiral because no mass term can be given to the anti-symmetric tensor field, and this field is an essential ingredient of the dynamics that generates the IR structure.

At the cubic level, the IR superpotential of the model contains two parts. One is identical to the ISS superpotential, with some fields getting supersymmetry-breaking mass splittings at tree-level. We will refer to these fields as messengers, in the spirit of GMSB models [6]. The second part of the superpotential contains couplings of the remaining pseudo-moduli to the messengers. As is well known by now [7, 8], this results in a calculable, rising potential at large field VEVs, so that the pseudo-moduli are stabilized.

It is not apriori clear however that the remaining pseudo-moduli are stabilized at the origin. The IR theory is an O’Raifeartaigh model, with all charges being 0 or 2. Such models were argued to preserve R-symmetry [9], and indeed, the mass-squared of the supersymmetry-breaking pseudo-modulus is positive at the origin. However, this conclusion applies only if one neglects cubic interactions that do not involve the supersymmetry breaking modulus¹. These interactions induce masses-squared for the remaining pseudo-moduli at one-loop, much like matter-messenger couplings in GMSB models, which are notorious for generating negative contributions to the scalar masses-squared at one-loop [11–13]². Still, in

¹ This loophole was recently pointed out in [10] too.

² The one-loop contributions of matter-messenger couplings to scalar masses were mainly calculated in

the present example, the resulting masses-squared are positive, so that the pseudo-moduli are stabilized at the origin with the R-symmetry unbroken.

It is worth noting that the IR superpotential we study involves two uncalculable couplings, and, as a result, the masses-squared of some of the pseudo-moduli at the origin are given by a non-trivial function of the ratio of these couplings. Surprisingly however, it turns out that the masses-squared are positive for any value of this ratio. This seems to hint at some general argument for the sign of the masses, and it would be interesting to understand this point.

Finally, a nice feature of the model is that one does not necessarily have to introduce a small mass scale by hand in order to obtain a calculable meta-stable minimum. Instead of adding a mass term in the UV, it is possible to add a higher-dimension term which becomes a linear term in the IR, with a coefficient that is naturally small.

We describe the theory we consider in Section II. In Section III we study a mass deformation. We comment on a non-renormalizable deformation in Section IV. Some details of the calculation of the mass-squared are presented in the Appendix.

II. THE SU(6) THEORY

We consider an SU(6) gauge theory with a single 3-index anti-symmetric tensor \mathcal{A} , and four flavors (see Table). The theory was shown in [4] to s-confine. The IR theory consists of the gauge invariants listed in the Table,

	$SU(6)$	$SU(4)_L$	$SU(4)_R$	$U(1)_1$	$U(1)_2$	$U(1)_R$
\mathcal{A}	$\begin{smallmatrix} \square \\ \square \\ \square \end{smallmatrix}$	1	1	0	-4	-1
Q	\square	\square	1	1	3	1
\bar{Q}	$\bar{\square}$	1	$\bar{\square}$	-1	3	1
$M \sim (Q\bar{Q})$		\square	$\bar{\square}$	0	6	2
$\Phi \sim (Q\mathcal{A}^2\bar{Q})$		\square	$\bar{\square}$	0	-2	0
$\bar{V} \sim (\mathcal{A}Q^3)$		$\bar{\square}$	1	3	5	2
$V \sim (\mathcal{A}\bar{Q}^3)$		1	\square	-3	5	2
$\bar{B} \sim (\mathcal{A}^3Q^3)$		$\bar{\square}$	1	3	-3	0
$B \sim (\mathcal{A}^3\bar{Q}^3)$		1	\square	-3	-3	0
$\tilde{\mathcal{A}} \sim (\mathcal{A}^4)$		1	1	0	-16	4

(1)

with the superpotential [4],

$$W = \frac{1}{\Lambda^{11}} \left[M_{IA} B_I \bar{B}_A + \Phi_{IA} (V_I \bar{B}_A + B_I \bar{V}_A) \right] + W_{NR}. \quad (2)$$

Here $I, A = 1, \dots, 4$, Λ is the dynamical scale of the theory, and

$$W_{NR} = \frac{1}{\Lambda^{11}} \left[V M \bar{V} \tilde{\mathcal{A}} + M \Phi^3 + \tilde{\mathcal{A}} M^3 \Phi \right] \quad (3)$$

the limit of small supersymmetry breaking [11, 12]. With the messenger spectrum of Minimal Gauge Mediation they vanish to leading order in the supersymmetry breaking as explained in [14]. There are also some examples for which these contributions were calculated for large supersymmetry breaking [13], with either sign possible.

where we suppressed global symmetry indices.

Since we are interested in vacua near the origin, with all field VEVs much smaller than Λ , the non-renormalizable part of the superpotential, W_{NR} , can be neglected. In order to give mass to the field $\tilde{\mathcal{A}}$, which only appears in W_{NR} , we can either add a singlet S to the theory, with the superpotential coupling

$$\Delta W = \frac{1}{M_{UV}^2} S (\mathcal{A}^4), \quad (4)$$

where M_{UV} is the UV cutoff scale, or simply add the superpotential term,

$$\frac{1}{M_{UV}^5} (\mathcal{A}^4)^2. \quad (5)$$

Either one of these becomes a mass term for $\tilde{\mathcal{A}}$ in the IR.

Following ISS, we could in principle add the superpotential

$$W_0 = m_{0IA} Q_I \bar{Q}_A + \frac{1}{M_{UV}} \lambda_{\Phi IA} Q_I \mathcal{A}^2 \bar{Q}_A, \quad (6)$$

where λ_{Φ} is dimensionless. In the IR theory, this becomes

$$W_0 = -(\mu_M)_{IA} M_{IA} - (\mu_{\Phi})_{IA} \Phi_{IA}, \quad (7)$$

with $\mu_M \sim \Lambda m_0$ and $\mu_{\Phi} \sim \Lambda^2/M_{UV}$. With matrices μ_M of rank greater than 1 and/or μ_{Φ} of rank greater than 2, supersymmetry is broken, since the F term for M_{IA} is

$$B_I \bar{B}_A - \mu_{IA} \quad (8)$$

and $B_I \bar{B}_A$ has rank 0 or 1. Similarly the F term for Φ_{IA} is

$$B_I \bar{V}_A + V_I \bar{B}_A - \mu_{IA} \quad (9)$$

and $B_I \bar{V}_A + V_I \bar{B}_A$ has maximal rank 2.

In the following, we will consider perturbations that preserve the maximal possible global symmetry, and therefore take μ_M and μ_{Φ} to be proportional to the identity. From the point of view of the IR theory, only one combination of M and Φ appears linearly and triggers supersymmetry breaking. Using the results of ISS, it is easy to see what happens. The field that appears linearly couples to some combination(s) of B and V . One then finds an ISS-like model involving one combination of M , Φ , and some combination of B - V with a supersymmetry-breaking minimum at the origin, and with all scalars apart from the Goldstones getting mass either at tree-level or at one-loop. The orthogonal combinations of $M - \Phi$ and $B - V$ couple to these through the superpotential and are therefore stabilized as well [7]. Here we will focus on the M perturbation for simplicity. From the point of view of the microscopic theory however, adding a linear term in Φ has some aesthetic advantage, since, in order to have a calculable minimum we want μ_M and/or μ_{Φ} much smaller than Λ . This automatically holds for μ_{Φ} , since the Φ tadpole originates from a non-renormalizable term.

III. A LINEAR TERM IN M

Consider first adding just a linear term in M , with $m_{0IA} \propto \delta_{IA}$ and $\lambda_{\Phi IA} = 0$. Written in terms of canonically-normalized gauge invariants, the superpotential of the IR theory is

$$W = h\mu M_{IA} (-\delta_{IA} + B_I \bar{B}_A) + \lambda \Phi_{IA} (V_I \bar{B}_A + B_I \bar{V}_A) . \quad (10)$$

Here we used the fact that the Kähler potential of the IR theory starts as³,

$$K = \frac{1}{\alpha_M^2} M^\dagger M + \frac{1}{\alpha_\Phi^2} \Phi^\dagger \Phi + \frac{1}{\alpha_B^2} (B^\dagger B + \bar{B}^\dagger \bar{B}) + \frac{1}{\alpha_V^2} (V^\dagger V + \bar{V}^\dagger \bar{V}) + \dots , \quad (11)$$

where the α 's are non-calculable order-one coefficients, and rescaled the fields M , Φ , B , V , defining

$$\mu = \frac{1}{\alpha_B^2} m_0 \Lambda , \quad h = \alpha_M \alpha_B^2 , \quad \lambda = \alpha_\Phi \alpha_B \alpha_V . \quad (12)$$

The model consists of two copies of the ISS fields, (M, B, \bar{B}) and (Φ, V, \bar{V}) with the same charges under all the global symmetries apart from $U(1)_2$. The first two terms of Eq. (10) are precisely the superpotential of the ISS SQCD model with three colors and four flavors, which has a minimum near the origin with all scalars (apart from the Goldstone bosons) in (M, B, \bar{B}) getting mass either at tree-level or at one-loop.

What happens to the remaining fields (Φ, V, \bar{V}) at this extremum? To answer this question, let us review the details of the ISS minimum. At this minimum, $M = 0$, and we can choose $B_4 = \bar{B}_4 = \mu$ up to global symmetries. Thus, the F term equation for M_{44} is satisfied. This is the maximal number of F -equations that can be satisfied in this case. Using small Latin indices for the unbroken $SU(3)_D$, the F terms of M_{ii} with $i = 1, 2, 3$ are nonzero, with $F_{M_{ii}} = -\mu^2$. The unbroken symmetry at the minimum is $SU(3)_D \times U(1)' \times U(1)_R$ where $U(1)'$ is a combination of $SU(4)_D$ and $U(1)_1$.

It will be convenient for our purposes to describe the spectrum of the ISS fields in analogy with minimal gauge mediation (MGM) models [6], splitting the fields according to their $SU(3)_D$ representations. We thus define, following ISS,

$$\begin{aligned} B_i &= b_i , & \bar{B}_i &= \bar{b}_i \\ B_4 &= \mu + B_4^+ + B_4^- , & \bar{B}_4 &= \mu + B_4^+ - B_4^- , \\ M_{ia} &= X_{ia} \\ M_{4a} &= z_a , & M_{i4} &= \bar{z}_i . \end{aligned} \quad (13)$$

Expanding around the minimum one has, from the first part of Eq. (10),

$$W = h\bar{b}Xb + h\mu(\bar{b}z + \bar{z}b) + h(\bar{z}b - \bar{b}z)B_4^- + h\mu M_{44}B_4^+ + \dots \quad (14)$$

where we omitted cubic terms involving M_{44} , B_4^\pm . The fields M_{44} and B_4^+ get mass at tree level. The $SU(3)_D$ singlet $X \equiv \text{Tr} M$ plays the role of the MGM singlet, with $F_X = h^2 \mu^2$, and

³ This form follows from the $SU(4)^2 \times U(1)^2 \times U(1)_R$ global symmetry, plus the exchange symmetry $Q \leftrightarrow \bar{Q}$, $\mathcal{A} \leftrightarrow \mathcal{A}^{dual}$, with $\mathcal{A}^{dual i_1 i_2 i_3} \equiv \epsilon^{i_1 \dots i_6} \mathcal{A}_{i_4 i_5 i_6}$ and with the vector superfield changing sign. This exchange symmetry also guarantees the equality of the coefficients of the last two terms of Eq. (10).

splits the masses-squared of the scalars in the messenger fields b, \bar{b} , while the supersymmetric masses of b and \bar{b} arise from their couplings to z, \bar{z} . All in all, the $b - z$ sector contains two fermions of masses $\pm h\mu$, two scalars of the same masses (from z and \bar{z}), and two scalars (from b, \bar{b}) with masses-squared $h^2\mu^2 \pm F_X = 0, 2h^2\mu^2$. The remaining pseudo-moduli X and B_4^- obtain masses at one-loop, through their superpotential couplings to the messengers.

Let us turn now to the fields Φ, V, \bar{V} . Splitting these according to their $SU(3)_D$ representations as in Eq. (13) we write

$$\begin{aligned} V_i &= v_i, & \bar{V}_i &= \bar{v}_i \\ V_4 &= V_4^+ + V_4^-, & \bar{V}_4 &= V_4^+ - V_4^-, \\ \Phi_{ia} &= Y_{ia} \\ \Phi_{4a} &= z'_a, & \Phi_{i4} &= \bar{z}'_i. \end{aligned} \tag{15}$$

so that the remaining piece of the superpotential Eq. (10) takes the form

$$W = \lambda\mu (\bar{z}'v + \bar{v}z') + \lambda\mu \Phi_{44}V_4^+ + \lambda Y (\bar{b}v + \bar{v}b) + \lambda V_4^- (\bar{b}z' - \bar{z}'b) + \dots \tag{16}$$

where we again neglected irrelevant cubic terms. At tree-level, just as in the ISS sector, the fields Φ_{44}, V_4^+ and $v - \bar{v}, z' - \bar{z}'$ get mass $\lambda\mu$ with Y, V_4^- remaining massless. These pseudo-moduli couple to the messengers through the superpotential Eq. (16) and are therefore stabilized [7, 8]. The reason is that, far from the origin (but for VEVs still smaller than Λ , where the theory is calculable) the potential for these fields can be reliably computed from the wave-function renormalizations of the light fields, as in [14]. This always results in a rising potential.

To find the masses near the origin, one must compute the Coleman-Weinberg potential. We present the result for arbitrary F_X in the Appendix. For $F_X = h^2\mu^2$ one finds,

$$m_Y^2 = \frac{1}{3}m_{V_4^-}^2 = \frac{1}{16\pi^2} \frac{\lambda^2 h^2}{(2 - \lambda^2/h^2)(1 - \lambda^2/h^2)} \left[\frac{\lambda^2}{h^2} \ln \left(\frac{\lambda^2}{h^2} \right) + 2 \left(1 - \frac{\lambda^2}{h^2} \right) \ln 2 \right] \mu^2, \tag{17}$$

which is positive for all values of the ratio λ^2/h^2 , so that these fields are stabilized at the origin.

With the addition of the linear term in M , the IR theory has an R-symmetry under which M has charge 2 so that B and \bar{B} can be chosen to have R-charge 0⁴. Therefore, the charges of V (\bar{V}) and Φ must sum to 2, and we can choose Φ to have R-charge 2 and V (\bar{V}) to have R-charge zero or vice versa. A priori, the cubic superpotential couplings of the pseudo-moduli to the messengers could have generated negative masses-squared for Φ (or V_4^-), leading to R-symmetry breaking. As we saw above, this is not the case. In fact, the result is positive even if we allow an arbitrary $F_X \neq h^2\mu^2$ in Eq. (16).

In addition to the $SU(4)_D \times U(1)_1 \times U(1)_R$ symmetry of the UV theory, the superpotential Eq. (10) preserves a global U(1), with, for example, V, \bar{V} having charge 1 and Φ having charge -1 . The only terms with $d \leq 3$ (in the IR fields) consistent with the IR symmetry are those already appearing in Eq. (10), but the singlet and adjoint pieces in M and Φ could now appear with different coefficients. Thus for example, one could add the term $\text{Tr} \Phi B \bar{V}$,

⁴ This symmetry happens to coincide with the anomaly-free $U(1)_R$ symmetry of the UV theory as given in the Table, but for this discussion all we care about is the effective R-symmetry of the IR theory.

or, (just as in the ISS model) $M_{adj}B\bar{B}$, where M_{adj} is the $SU(4)$ adjoint part of M . The superpotential of the IR model is therefore not generic. However, if these terms are added by hand in the UV theory, they are suppressed by the UV cutoff scale, and their contributions to the masses are therefore smaller than the radiatively-generated contributions.

IV. A LINEAR TERM IN Φ

Let us briefly comment on adding a linear term in Φ only, setting $m_0 = 0$ and $\lambda_\phi \neq 0$ in Eq. (6). It is convenient to define the combinations

$$p_I \propto V_I + B_I, \quad q_I \propto V_I - B_I, \quad (18)$$

and similarly for the barred fields. The superpotential then takes the form

$$W = \lambda\Phi_{IA} (p_I\bar{p}_A + q_I\bar{q}_A - \mu_\Phi\delta_{IA}) + \lambda M_{IA} (p_I\bar{p}_A + q_I\bar{q}_A + p_I\bar{q}_A - q_I\bar{p}_A), \quad (19)$$

and the supersymmetry breaking scale, $\mu_\Phi \sim \Lambda^2/M_{UV}$ is naturally small. In this case, two of the F -term equations for Φ can be solved, with two entries of p, q getting VEVs. The pseudo-moduli will again be stabilized at one-loop by the superpotential couplings to the messengers. It would be interesting to study the fate of the R symmetry in models that contain both the $\text{Tr}M$ and the $\text{Tr}\Phi$ perturbations.

V. CONCLUSIONS

As Intriligator, Seiberg, and Shih [2] demonstrated, meta-stable supersymmetry-breaking vacua appear in the simplest theories. This suggests that local supersymmetry breaking vacua may occur quite generically, in both chiral and non-chiral theories. In this note, we discussed generalization of the ISS vacua to chiral theories in s-confining examples, which are particularly easy to analyze. It would be interesting to go beyond s-confinement, and to explore chiral theories that have weakly coupled IR duals in the search for local supersymmetry breaking vacua, whether they are obtained by rank-condition breaking or by other means. Even chiral theories with global, supersymmetry breaking vacua, might possess additional, local minima with novel features.

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Appendix A: The Coleman-Weinberg calculation

It is instructive to calculate the 1-loop masses for arbitrary F . We can write the one-loop correction to the vacuum energy as

$$V_{1-loop} = \frac{1}{64\pi^2} \left(\text{Tr} m_B^4 \log \frac{m_B^2}{\Lambda^2} - \text{Tr} \bar{m}_B^4 \log \frac{\bar{m}_B^2}{\Lambda^2} \right) \quad (\text{A1})$$

where m_B denotes the boson masses, and \bar{m}_B denotes the same mass in the supersymmetric limit ($F = 0$). To obtain this result we used the fact that the correction vanishes for $F = 0$, and that the fermion masses are F -independent.

To derive the Y mass we write $m_B^2 = a + b|Y|^2$ neglecting higher orders in Y . The vanishing of the supertrace gives

$$\Sigma_i a_i = \Sigma_i \bar{a}_i \quad (\text{A2})$$

$$\Sigma_i b_i = \Sigma_i \bar{b}_i \quad (\text{A3})$$

$$\Sigma_i a_i b_i = \Sigma_i \bar{a}_i \bar{b}_i \quad (\text{A4})$$

where again, bars denote the supersymmetric quantities. The last equality implies that V_{1-loop} is finite. The Y mass-squared is then given by

$$m_Y^2 = \frac{1}{32\pi^2} \Sigma_i \left(a_i b_i \log a_i - \bar{a}_i \bar{b}_i \log \bar{a}_i \right). \quad (\text{A5})$$

One then finds

$$m_Y^2 = \frac{1}{16\pi^2} \frac{h^2 \lambda^2 \mu^2}{(1-r-f)(1-r+f)(1-r)} \left[-2r^2 \ln(r) f^2 + (1-r) \left((1+f)^2 (1-r-f) \ln(1+f) + (1-f)^2 (1-r+f) \ln(1-f) \right) \right] \quad (\text{A6})$$

with $r = \lambda^2/h^2$, $f = F/(h^2\mu^2)$, which reduces to Eq. (17) for $f = 1$. We note that Eq. (A6) is positive for all values of $f \leq 1$.

We also note that, unlike the one-loop contributions of matter-messenger couplings in MGM models, this contribution does not vanish at $O(F^2)$, because the supersymmetric mass of the messengers does not arise from X .

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